

# Analysis of direct CP violation in $B^- \rightarrow D^0 D_s^-, D^0 D^-$ decays

A. K. Giri<sup>1</sup>, R. Mohanta<sup>2</sup> and M. P. Khanna<sup>1</sup>

<sup>1</sup> *Physics Department, Panjab University,  
Chandigarh - 160 014, India*

<sup>2</sup> *School of Physics, University of Hyderabad,  
Gachibowli, Hyderabad - 500 046, India*

February 1, 2008

## Abstract

We investigate the possibility of observing the direct CP violation in the decay modes  $B^- \rightarrow D^0 D_s^-$  and  $D^0 D^-$  within the Standard Model. Including the contributions arising from the tree, annihilation, QCD as well as electroweak penguins with both time- and space-like components, we find that the direct CP asymmetry in  $B^- \rightarrow D^0 D_s^-$  is very small  $\sim 0.2\%$  but in  $B^- \rightarrow D^0 D^-$  decay it can be as large as 4%. Approximately  $10^7$  charged  $B$  mesons are required to experimentally observe the CP asymmetry parameter for the later case. Since this is easily accessible with the currently running B factories, the decay mode  $B^- \rightarrow D^0 D^-$  may be pursued to look for CP violation.

PACS Numbers. : 11.30 Er, 13.25 Hw

## 1 Introduction

CP violation is one of the least understood phenomena in particle physics [1, 2, 3], although it was observed in  $K^0 - \bar{K}^0$  mixing system more than 35 years ago. In the standard model (SM), CP violation arises from a complex phase in the Cabibbo-Kobayashi-Maskawa quark mixing matrix [4]. Outside the Kaon system, decays of  $B$  mesons provide rich ground for investigating CP violation [5, 6]. Within the SM, the CP violation is often characterized by the so-called unitarity triangle [7]. By measuring CP violating rate asymmetries in  $B$  decays, one can extract  $\alpha$ ,  $\beta$  and  $\gamma$ , the three interior angles of the unitarity triangle. The sum of these three angles must be equal to  $180^\circ$  in the SM with three generations. At present we are at the beginning of the  $B$ -factory era

in particle physics, which will provide us valuable insights to understand the phenomena of CP violation. One of the main programmes of the presently running and the upcoming  $B$  factories is to measure the size of CP violation in as many  $B$  decay modes as possible so as to establish the pattern of CP violation in various  $B$  decays. Among the most interesting  $B$  decay channels, the “gold plated” mode  $B_d \rightarrow J/\psi K_s$ , [8] allows the determination of the angle  $\beta$  of the unitarity triangle of CKM matrix. Recent measurement of CP asymmetry in the  $B^0 \rightarrow J/\psi K^0$  and other related processes e.g.  $\psi' K^0$ ,  $\eta_c K^0$  etc by the BELLE [9] and BaBar [10] detectors at the KEK and SLAC  $B$  factories together with the earlier measurement of CDF [11] constitute the first significant signal of CP violation outside the neutral Kaon system.

While the most promising proposal for observing CP violation in the  $B$ -system involves the mixing between neutral  $B$  mesons [1], the decays of charged  $B$  mesons are also of particular importance for establishing the detailed nature of CP violation. Since charged  $B$  mesons cannot mix, a measurement of the CP violating observable in these decays would be a clear sign of “direct CP violation” which has been searched for in  $K$ -system for quite long with indefinite success. Only recently, such kind of CP violating effect has been observed in the  $K$  system by the NA48 [12] and KTeV [13] collaborations. For the bottom meson case usually the charmless rare  $B$  decay modes are preferred to study the direct CP violation as these decay modes proceed with more than one Feynman diagrams. In this paper we would like to look for some additional decay channels which could help us in establishing the presence of CP violation as quickly as possible. For this purpose we investigate the direct CP violating effects in the decays of charged  $B$  mesons to two charmed mesons i.e.  $B^- \rightarrow D^0 D_s^-$  and  $D^0 D^-$ . It is worth emphasizing that these decay modes are flavor self tagging processes which should be favored for experimental reconstructions. The decay mode  $B^- \rightarrow D^0 D_s^-$  has already been observed experimentally with a branching ratio  $(1.3 \pm 0.4)\%$  and the upper limit for  $B^- \rightarrow D^0 D^-$  channel is found to be  $< 6.7 \times 10^{-3}$  [14]. These decay modes which are described by the quark level transitions as  $b \rightarrow c\bar{c}q$  ( $q = s/d$  for  $D_s^-/D^-$  in the final state) proceed through three distinct type of flavor topologies. These are : the color allowed but Cabibbo suppressed tree, annihilation and the QCD as well as electroweak penguin diagrams. To get significant direct CP violation one would require two interfering amplitudes of comparable strengths, with different strong and weak phases. The weak phases arise from the superposition of various penguin contributions and the usual tree diagrams. The strong phases are generated by the perturbative penguin loops (hard final state interaction) [15] or final state interactions involving two different isospins. Since the decay modes we considered here are single isospin channels i.e. the final states  $D^0 D_s^-$  and  $D^0 D^-$  are with isospin  $I = 1/2$  and 1 respectively the second type of FSI strong phase differences are absent for these channels. Therefore at the first sight it appears that direct CP violating effects in these channels would be negligibly small as

the tree contribution dominates over the other diagrams and thus have been overlooked in the literature. But detailed calculation shows that it is indeed not so. In fact the CP violating effects in  $B^- \rightarrow D^0 D^-$  channel can be as large as few percent level which can be experimentally accessible in the first round of  $B$  factories. The reason for the existence of such a significant CP violating parameter may be due to the fact that the tree diagram for  $b \rightarrow c\bar{c}d$  transition is although colour allowed, it is doubly Cabibbo suppressed, and hence its magnitude is not very much larger than the penguin contributions. CP violating effects in the decays of neutral  $B$  meson into double charmed mesons have been extensively studied in Refs. [16-19], where it has been shown that, these channels can be used as an alternative method to the  $J/\psi K_s$  mode for the extraction of the angle  $\beta$ .

In our analysis, we use the standard theoretical framework to study the nonleptonic  $B^- \rightarrow D^0 D_s^-(D^-)$  decay modes, which is based on the effective Hamiltonian approach in conjunction with the factorization hypothesis. The short distance QCD corrected Hamiltonian is calculated to next-to-leading order. The renormalization scheme and scale problems with factorization approach for matrix elements can be circumvented by employing the scale and scheme independent effective Wilson coefficients. In the literature the contributions of space-like penguins are neglected assuming form factor suppression. But as pointed out in Ref. [20] the effect of space like penguin amplitudes can be remarkably enhanced by the hadronic matrix elements involving  $(V - A)(V + A)$  or  $(S + P)(S - P)$  currents. Therefore we have included the space and time like contributions of both QCD and EW penguins, the annihilation contribution in addition to the dominant tree diagrams. Assuming the factorization approximation, the matrix elements of the tree and time-like penguin diagrams have been calculated in the BSW model [21], whereas for the evaluation of the matrix elements of the space and annihilation diagrams we have employed the Lepage and Brodsky model [22].

The paper is organized as follows. In section II we briefly discuss the effective Hamiltonian together with the quark level matrix elements and the numerical value of the Wilson coefficients in the effective Hamiltonian approach. Assuming the factorization approximation, the matrix elements of tree and time-like penguins are evaluated in the BSW model and for the space-like and annihilation diagrams we use LB (Lepage and Brodsky) model. Determination of the CP violating asymmetry is presented in section III and section IV contains our conclusion.

## 2 Framework

The effective Hamiltonian  $\mathcal{H}_{eff}$  for the decay modes  $B^- \rightarrow D^0 D_s^-$  and  $D^0 D^-$  which are described by the quark level transitions  $b \rightarrow c\bar{c}q$  (where  $q = s$  for the former and  $d$  for the later) have three classes of flavour topologies : the

dominant tree, annihilation, and both QCD as well as electroweak penguins given by [5]

$$\begin{aligned} \mathcal{H}_{eff} = & \frac{G_F}{\sqrt{2}} \left\{ \lambda_u [c_1(\mu) O_1^u(\mu) + c_2(\mu) O_2^u(\mu)] + \lambda_c [c_1(\mu) O_1^c(\mu) + c_2(\mu) O_2^c(\mu)] \right. \\ & \left. + (\lambda_u + \lambda_c) \sum_{i=3}^{10} c_i(\mu) O_i(\mu) \right\} + \text{h.c.} , \end{aligned} \quad (1)$$

where  $\lambda_u = V_{ub}V_{uq}^*$  and  $\lambda_c = V_{cb}V_{cq}^*$  and  $c_i(\mu)$  are the Wilson coefficients evaluated at the renormalization scale  $\mu$ . The four fermion operators  $O_{1-10}$  are given as

$$\begin{aligned} O_1^u &= (\bar{u}b)_{V-A}(\bar{q}u)_{V-A} , & O_2^u &= (\bar{u}_\alpha b_\beta)_{V-A}(\bar{q}_\beta u_\alpha)_{V-A} , \\ O_1^c &= (\bar{c}b)_{V-A}(\bar{q}c)_{V-A} , & O_2^c &= (\bar{c}_\alpha b_\beta)_{V-A}(\bar{q}_\beta c_\alpha)_{V-A} , \\ O_{3(5)} &= (\bar{q}b)_{V-A} \sum_{q'} (\bar{q}'q')_{V-A(V+A)} , \\ O_{4(6)} &= (\bar{q}_\alpha b_\beta)_{V-A} \sum_{q'} (\bar{q}'_\beta q'_\alpha)_{V-A(V+A)} , \\ O_{7(9)} &= \frac{3}{2} (\bar{q}b)_{V-A} \sum_{q'} e_{q'} (\bar{q}'q')_{V+A(V-A)} , \\ O_{8(10)} &= \frac{3}{2} (\bar{q}_\alpha b_\beta)_{V-A} \sum_{q'} e_{q'} (\bar{q}'_\beta q'_\alpha)_{V+A(V-A)} , \end{aligned} \quad (2)$$

where  $O_{1,2}$  are the tree level current-current operators,  $O_{3-6}$  the QCD penguin operators and  $O_{7-10}$  the EW penguin operators.  $(\bar{q}_1 q_2)_{(V \pm A)}$  denote the usual  $(V \pm A)$  currents. The sum over  $q'$  runs over the quark fields that are active at the scale  $\mu = O(m_b)$  i.e.  $(q' \in u, d, s, c, b)$ . The Wilson coefficients depend (in general) on the renormalization scheme and the scale  $\mu$  at which they are evaluated. In the next to leading order their values obtained in the naive dimensional regularization (NDR) scheme at  $\mu = m_b(m_b)$  as [23]  $c_1 = 1.082$ ,  $c_2 = -0.185$ ,  $c_3 = 0.014$ ,  $c_4 = -0.035$ ,  $c_5 = 0.009$ ,  $c_6 = -0.041$ ,  $c_{7/\alpha} = -0.002$ ,  $c_{8/\alpha} = 0.054$ ,  $c_{9/\alpha} = -1.292$  and  $c_{10/\alpha} = 0.263$ .

However, the physical matrix elements  $\langle P_1 P_2 | \mathcal{H}_{eff} | B \rangle$  are obviously independent of both the scheme and the scale. Hence the dependence on the Wilson coefficients must be compensated by a comensurate calculation of the hadronic matrix elements in a nonperturbative framework such as lattice QCD. Presently, this is not a viable strategy as the calculation of the matrix elements  $\langle P_1 P_2 | O_i | B \rangle$  is beyond the scope of the current lattice technology. However, perturbation theory comes to (partial) rescue; with the help of which one-loop matrix elements can be rewritten in terms of the operators and the effective Wilson coefficients  $c_i^{eff}$  which are scheme and scale independent :

$$\langle qq' \bar{q}' | \mathcal{H}_{eff} | b \rangle = \sum_{i,j} c_i^{eff}(\mu) \langle qq' \bar{q}' | O_j | b \rangle^{tree} . \quad (3)$$

The effective Wilson coefficients  $c_i^{eff}(\mu)$  may be expressed as [24]

$$\begin{aligned}
c_1^{eff}|_{\mu=m_b} &= c_1(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{1i} c_i(\mu) , \\
c_2^{eff}|_{\mu=m_b} &= c_2(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{2i} c_i(\mu) , \\
c_3^{eff}|_{\mu=m_b} &= c_3(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{3i} c_i(\mu) - \frac{\alpha_s}{24\pi} (C_t + C_p + C_g) , \\
c_4^{eff}|_{\mu=m_b} &= c_4(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{4i} c_i(\mu) + \frac{\alpha_s}{8\pi} (C_t + C_p + C_g) , \\
c_5^{eff}|_{\mu=m_b} &= c_5(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{5i} c_i(\mu) - \frac{\alpha_s}{24\pi} (C_t + C_p + C_g) , \\
c_6^{eff}|_{\mu=m_b} &= c_6(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{6i} c_i(\mu) + \frac{\alpha_s}{8\pi} (C_t + C_p + C_g) , \\
c_7^{eff}|_{\mu=m_b} &= c_7(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{7i} c_i(\mu) + \frac{\alpha}{8\pi} C_e , \\
c_8^{eff}|_{\mu=m_b} &= c_8(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{8i} c_i(\mu) , \\
c_9^{eff}|_{\mu=m_b} &= c_9(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{9i} c_i(\mu) + \frac{\alpha}{8\pi} C_e , \\
c_{10}^{eff}|_{\mu=m_b} &= c_{10}(\mu) + \frac{\alpha_s}{4\pi} \left( \gamma^{(0)T} \ln \frac{m_b}{\mu} + \hat{r}^T \right)_{10i} c_i(\mu) . \tag{4}
\end{aligned}$$

where  $\hat{r}^T$  and  $\gamma^{(0)T}$ , the transpose of the matrices  $\hat{r}$  and  $\gamma^{(0)}$ , arise from the vertex corrections to the operators  $O_1 - O_{10}$  derived in [25], which are explicitly given in Ref. [26]

The quantities  $C_t$ ,  $C_p$ ,  $C_e$  and  $C_g$  are arising from the penguin type diagrams of the operators  $O_{1,2}$ , the QCD penguin type diagrams of the operators  $O_3 - O_6$ , the electroweak penguin type diagrams of  $O_{1,2}$  and the tree level diagrams of the dipole operator  $O_g$  respectively, which are given in the NDR scheme (after  $\overline{\text{MS}}$  renormalization) by

$$\begin{aligned}
C_t &= - \left( \frac{\lambda_u}{\lambda_t} \tilde{G}(m_u) + \frac{\lambda_c}{\lambda_t} \tilde{G}(m_c) \right) c_1 \\
C_p &= [\tilde{G}(m_s) + \tilde{G}(m_b)] c_3 + \sum_{i=u,d,s,c,b} \tilde{G}(m_i) (c_4 + c_6) \\
C_g &= - \frac{2m_b}{\sqrt{\langle k^2 \rangle}} c_g^{eff} , \quad c_g^{eff} = -1.043
\end{aligned}$$

$$C_e = -\frac{8}{9} \left( \frac{\lambda_u}{\lambda_t} \tilde{G}(m_u) + \frac{\lambda_c}{\lambda_t} \tilde{G}(m_c) \right) (c_1 + 3c_2)$$

$$\tilde{G}(m_q) = \frac{2}{3} - G(m_q, k, \mu) \quad (5)$$

$$G(m, k, \mu) = -4 \int_0^1 dx \, x(1-x) \ln \left( \frac{m^2 - k^2 x(1-x)}{\mu^2} \right), \quad (6)$$

It should be noted that the quantities  $C_t$ ,  $C_p$ ,  $C_e$  and  $C_g$  depend on the CKM matrix elements, the quark masses, the scale  $\mu$  and  $k^2$ , the momentum transferred by the virtual particles appearing in the penguin diagrams. In the factorization approximation there is no model independent way to keep track of the  $k^2$  dependence; the actual value of  $k^2$  is model dependent. From simple kinematics [27] one expects  $k^2$  to be typically in the range

$$\frac{m_b^2}{4} \leq k^2 \leq \frac{m_b^2}{2}. \quad (7)$$

Since the branching ratio and the CP asymmetry depend crucially on the parameter  $k^2$ , here we would like to take a specific value for it based on valence quark approximation instead of the conventionally used value  $k^2 = m_b^2/2$ . As discussed in Ref. [20] the averaged value of the squared momentum transfer for  $B^-(b\bar{u}) \rightarrow D^0(c\bar{u})D_q^-(q\bar{c})$  is given as

$$\langle k^2 \rangle = m_b^2 + m_q^2 - 2m_b E_q \quad (8)$$

where the energy of the quark  $q$  in the final  $D_q^-$  particle is determinable from

$$E_q + \sqrt{E_q^2 - m_q^2 + m_c^2} + \sqrt{4E_q^2 - 4m_q^2 + m_c^2} = m_b \quad (9)$$

for time-like penguin channels; or from

$$E_q + \sqrt{E_q^2 - m_q^2 + m_u^2} = m_b + m_u \quad (10)$$

for space-like penguin diagrams.  $m_b$ ,  $m_q$  and  $m_c$  denote the masses of the decaying  $b$ -quark, daughter  $q$ -quark and the  $c$ -quark (created as  $c\bar{c}$  pair from the virtual gluon, photon or  $Z$ -particle in the penguin loop). For numerical calculations, we have taken the CKM matrix elements expressed in terms of the Wolfenstein parameters with values  $A = 0.815$ ,  $\lambda = \sin \theta_c = 0.2205$ ,  $\rho = 0.175$  and  $\eta = 0.37$  [26]. The choice of  $\rho$  and  $\eta$  correspond to the CKM triangle:  $\alpha = 91^\circ$ ,  $\beta = 24^\circ$  and  $\gamma = 65^\circ$ . At scale  $\mu \sim m_b$ , we use the current quark masses as [26]  $m_u(m_b) = 3.2$  MeV,  $m_d(m_b) = 6.4$  MeV,  $m_s(m_b) = 90$  MeV,  $m_c(m_b) = 0.95$  GeV and  $m_b(m_b) = 4.34$  GeV. With the specific value of  $k^2$  obtained from Eqns. (8-10), we obtain the values of the effective renormalization scheme and scale independent Wilson coefficients for  $b \rightarrow s$  and  $b \rightarrow d$  transitions as given in Table-1.

Now we want to calculate the matrix element  $\langle D_q^- D^0 | O_i | B^- \rangle$  using the factorization approximation, where  $O_i$  are the four quark current operators listed above. In this approximation, the hadronic matrix elements of the four quark operators  $(\bar{c}b)_{(V-A)}(\bar{q}c)_{(V-A)}$  split into the product of two matrix elements,  $\langle D^0 | (\bar{c}b)_{(V-A)} | B^- \rangle$  and  $\langle D_q^- | (\bar{q}c)_{(V-A)} | 0 \rangle$  where Fierz transformation has been used so that flavor quantum numbers of the currents match with those of the hadrons. Since Fierz rearranging yields operators which are in the color singlet-singlet and octet-octet forms, this procedure results, in general, in matrix elements which have the right flavor quantum numbers but involve both singlet-singlet and octet-octet current operators. However, there is no experimental information available for the octet-octet part. So in the factorization approximation, one discards the color octet-octet piece and compensates this by treating  $N_c$ , the numbers of colors as a free parameter, and its value is extracted from the data of two body nonleptonic decays.

The matrix elements of the  $(V-A)(V+A)$  operators i.e. ( $O_6$  &  $O_8$ ) can be transformed into  $(V-A)(V-A)$  form by using Fierz ordering and the Dirac equation, which are given as

$$\langle D_q^- D^0 | O_6 | B^- \rangle = R_q \langle D_q^- D^0 | O_4 | B^- \rangle \quad (11)$$

with

$$R_q = \frac{2m_{D_q^-}}{(m_b - m_c)(m_q + m_c)}, \quad (12)$$

where the quark masses are the current quark masses. The same relation works for  $O_8$ .

Hence, one obtains the transition amplitude for  $B^- \rightarrow D_s^- D^0$  and  $D^- D^0$  as (where the factor  $G_F/\sqrt{2}$  is suppressed)

$$\begin{aligned} A(B^- \rightarrow D_s^- D^0) &= \lambda_u \left\{ \left( a_4 + a_{10} + (a_6 + a_8)R_s \right) X^{(BD^0, D_s^-)} \right. \\ &+ \left. \left( a_1 + a_4 + a_{10} + (a_6 + a_8)R'_s \right) X^{(B, D^0 D_s^-)} \right\} \\ &+ \lambda_c \left\{ \left( a_1 + a_4 + a_{10} + (a_6 + a_8)R_s \right) X^{(BD^0, D_s^-)} \right. \\ &+ \left. \left( a_4 + a_{10} + (a_6 + a_8)R'_s \right) X^{(B, D^0 D_s^-)} \right\} \end{aligned} \quad (13)$$

$$\begin{aligned} A(B^- \rightarrow D^- \bar{D}^0) &= \lambda_u \left\{ \left( a_4 + a_{10} + (a_6 + a_8)R_d \right) X^{(BD^0, D^-)} \right. \\ &+ \left. \left( a_1 + a_4 + a_{10} + (a_6 + a_8)R'_d \right) X^{(B, D^0 D^-)} \right\} \\ &+ \lambda_c \left\{ \left( a_1 + a_4 + a_{10} + (a_6 + a_8)R_d \right) X^{(BD^0, D^-)} \right. \\ &+ \left. \left( a_4 + a_{10} + (a_6 + a_8)R'_d \right) X^{(B, D^0 D^-)} \right\} \end{aligned} \quad (14)$$

where

$$\begin{aligned} X^{(BD^0, D_q^-)} &= \langle D_s^- | (\bar{q}c) | 0 \rangle \langle D^0 | (\bar{c}b) | B \rangle \\ X^{(B, D_q^- D^0)} &= \langle D^0 D_q^- | (\bar{q}c) | 0 \rangle \langle 0 | (\bar{u}b) | B \rangle . \end{aligned} \quad (15)$$

$X^{(BD^0, D_q^-)}$  denotes matrix elements of the tree and time like penguins where as  $X^{(B, D_q^- D^0)}$  stand for the annihilation and space-like amplitudes.

$$R_{q'} = \frac{2m_B^2}{(m_q - m_u)(m_b + m_u)} , \quad (16)$$

arises from the transformation of  $(V-A)(V+A)$  operators into  $(V-A)(V-A)$  form for space-like penguins. It should be noted that  $\lambda_u = V_{ub}V_{us}^*$  for  $B^- \rightarrow D^0 D_s^-$  whereas  $\lambda_u = V_{ub}V_{ud}^*$  for  $B^- \rightarrow D^0 D^-$  and similar expressions for  $\lambda_c$ .

The coefficients  $a_1, a_2 \dots a_{10}$  are combinations of the effective Wilson coefficients given as

$$a_{2i-1} = c_{2i-1}^{eff} + \frac{1}{N_c^{eff}} c_{2i}^{eff} \quad a_{2i} = c_{2i}^{eff} + \frac{1}{N_c^{eff}} c_{2i-1}^{eff} \quad i = 1, 2 \dots 5 , \quad (17)$$

where  $N_c^{eff}$  is the effective number of colors treated as free parameter in order to model the nonfactorizable contributions to the matrix elements and its value can be extracted from the two body nonleptonic  $B$  decays. A recent analysis of  $B \rightarrow D\pi$  data gives  $N_c^{eff} \sim 2$  [28]. Therefore, in our analysis, we take two sets of values for  $N_c^{eff}$  i.e.,  $N_c^{eff} = 2$  and  $N_c^{eff} = 3$ , which characterizes naive factorization.

The factorized hadronic matrix elements are evaluated using the BSW model [21], which are given as

$$X^{(BD^0, D_q^-)} = if_{D_q} F_0^{BD}(m_{D_q}^2)(m_B^2 - m_{D^0}^2) \quad (18)$$

The matrix element of the annihilation and space-like penguins are given as [20]

$$\langle D^0 D_q^- | (\bar{q}u)(\bar{u}b) | B^- \rangle = if_B f_+^a(m_B^2) \left[ m_{D_q}^2 - m_{D^0}^2 - \frac{m_{D_q} - m_{D^0}}{m_{D_q} + m_{D^0}} m_B^2 \right] , \quad (19)$$

where the value of the annihilation form factor is given as  $f_+^a(m_B^2) = i16\pi\alpha_s f_B^2/m_B^2$  [22].

After obtaining the transition amplitude, the branching ratio is given as

$$BR = \frac{|\mathbf{p}|}{8\pi m_B^2} \frac{|A(B^- \rightarrow D^0 D_q^-)|^2}{\Gamma} , \quad (20)$$

where  $|\mathbf{p}|$  is the momentum of the emitted particles and  $\Gamma$  is the total decay width.



Using eqns (13)-(19) we obtain the transition amplitude (in the unit of  $G_F/\sqrt{2}$ ) as

$$A(B^- \rightarrow D^0 D_s^-) = \lambda_u(0.1898 - i0.6483) + \lambda_c(0.1889 + i4.418) \\ [\lambda_u(0.2019 - i0.6817) + \lambda_c(0.201 + i4.698)] , \quad (21)$$

$$A(B^- \rightarrow D^0 D^-) = \lambda_u(0.2259 - i0.5616) + \lambda_c(0.2259 + i4.8185) \\ [\lambda_u(0.2393 - i0.589) + \lambda_c(0.2393 + i5.124)] , \quad (22)$$

where we have used the decay constants (in MeV) as  $f_{D_s} = 280$ ,  $f_D = 300$  [14] and  $f_B=180$  [29]. In the above equations, the upper values correspond to  $N_c^{eff} = 2$  and the lower bracketed values to  $N_c^{eff} = 3$ .

### 3 CP Violating Asymmetry

For charge  $B^\mp$  decays, the CP violating rate asymmetries in partial decay rates are defined as

$$a_{cp} = \frac{\Gamma(B^- \rightarrow f^-) - \Gamma(B^+ \rightarrow f^+)}{\Gamma(B^- \rightarrow f^-) + \Gamma(B^+ \rightarrow f^+)} . \quad (23)$$

As these decays are all self tagging the measurement of these CP violating asymmetry is essentially a counting experiment in well defined final states. Their rate asymmetries require both weak and strong phase differences in interfering amplitudes. The weak phase difference arises from the superposition of amplitudes from various tree (current-current) and penguin diagrams. The strong phase which are needed to obtain nonzero values for  $a_{cp}$  are generated by absorptive parts in penguin diagrams (hard final state interactions).

For the  $B$  meson decaying to a final state  $f$  and the charge conjugated  $B^- \rightarrow f$  we may, without any loss of generality, write the transition amplitude as

$$A(f) = \lambda_u A_u e^{i\delta_u} + \lambda_c A_c e^{i\delta_c} \quad (24)$$

$$\bar{A}(f) = \lambda_u^* A_u e^{i\delta_u} + \lambda_c^* A_c e^{i\delta_c} \quad (25)$$

where  $\lambda_i = V_{ib}V_{iq}^*$ ,  $A_u$  and  $A_c$  denote the contribution from penguin operators proportional to the product of CKM matrix elements  $\lambda_u$  and  $\lambda_c$  respectively. The corresponding strong phases are denoted by  $\delta_u$  and  $\delta_c$  respectively.

Thus the direct CP violating asymmetry is given as,

$$a_{cp} = \frac{-2 \operatorname{Im}(\lambda_u \lambda_c^*) \operatorname{Im}(A_u A_c^*)}{|\lambda_u A_u|^2 + |\lambda_c A_c|^2 + 2 \operatorname{Re}(\lambda_u \lambda_c^*) \operatorname{Re}(A_u A_c^*)} \\ = \frac{2 \sin \gamma \sin(\delta_u - \delta_c)}{|\frac{\lambda_u A_u}{\lambda_c A_c}| + |\frac{\lambda_c A_c}{\lambda_u A_u}| + 2 \cos \gamma \cos(\delta_u - \delta_c)} , \quad (26)$$

where the weak phases entering in the  $b \rightarrow s/d$  transition is equal to  $(-\gamma)$ , as we are using Wolfenstein approximation in which  $\lambda_c$  has no weak phase and the phase of  $\lambda_u$  is  $-\gamma$ . The strong phase  $(\delta_u - \delta_c)$  is caused by the final state interactions.

The strong phases are given by,

$$\sin(\delta_u - \delta_c) = \frac{1}{|A_u A_c|} (\text{Im} A_u \text{Re} A_c - \text{Im} A_c \text{Re} A_u) \quad (27)$$

$$\cos(\delta_u - \delta_c) = \frac{1}{|A_c A_u|} (\text{Re} A_u \text{Re} A_c + \text{Im} A_u \text{Im} A_c) \quad (28)$$

## 4 Conclusion

Using the next-to-leading order QCD corrected effective Hamiltonian, the scale and scheme independent Wilson coefficients, we have systematically studied the two charm hadronic decay modes  $B^- \rightarrow D^0 D_s^-$  and  $D^0 D^-$  within the framework of generalized factorization. The nonfactorizable contributions are parametrized in terms of  $N_c^{eff}$ , the effective number of colors. For numerical calculations, we have used two different sets of values for this parameters: (i)  $N_c^{eff} = 2$ , (ii)  $N_c^{eff} = 3$  which holds for naive factorization. The existence of a direct CP violating rate asymmetry requires two interfering amplitudes having different CP nonconserving weak phases and CP conserving strong phases. The former may arise either from the Standard Model CKM matrix or from new physics while the latter may arise from the absorptive part of a penguin diagram or from final state interaction effects of two different isospins. Since the channels we considered here are single isospin channels, the second class of strong phase differences do not arise for these channels. In our analysis, the weak phases are due to the CKM matrix and the strong phase differences arise due to absorptive part of penguin diagrams. The branching ratio and the CP violating asymmetry parameter are estimated using Eqs.(20) and (26) and are presented in Table-2

From the results we have observed the following:

1. The predicted branching ratio for the decay mode  $B^- \rightarrow D^0 D_s^-$  agrees very well with the experimental value for  $N_c^{eff} = 2$ , and the CP violating parameter for this mode is quite small.
2. The branching ratio for the decay mode  $B^- \rightarrow D^0 D^-$  lies below the present experimental upper bound and the CP violating parameter for this mode is quite significant. The number of charged  $B$  mesons required to observe this CP violating signal to three standard deviation is given as  $N_B = 9/(BR \times a_{CP}^2) \approx 7.9 \times 10^6$ , which is easily accessible with the currently running  $B$  factories.

It has been emphasized in Refs. [16-19] that the neutral B meson decay modes to two charmed mesons can be used to measure the unitarity angle  $\beta$

as alternative to the gold plated mode  $B \rightarrow J/\psi K$ . We argue further here that the mode  $B^- \rightarrow D^0 D^-$  can be used to quickly settle down the search for observing direct CP violation outside Kaon system, if SM description of CP violation is correct or else could provide us a clear indication of the presence of new physics. It should be noted here that the decay mode is flavor self-tagging and hence experimentally favourable. Furthermore, since the branching ratio is  $O(10^{-3})$  in our case, which is larger than the other competitive mode like  $B \rightarrow \pi K$  where the branching ratio is  $O(10^{-5})$  [14] and the direct CP violation parameter  $a_{cp}$  is significantly large i.e., we have obtained  $a_{cp}$  around 4 % as to that of  $-1.4$  % in ref [30], this decay mode is certainly a better candidate to observe CP violation in the first round of B-factory experiments.

To summarize, since the modes we consider are direct decays and not time dependent, they may be observed in any experimental setting where large number of  $B$  mesons are produced. Apart from the SLAC and KEK asymmetric  $B$  factories these include CLEO and hadronic  $B$  experiments such as HERA-b, BTeV, Collider Detectors at Fermilab (CDF), D0 and CERN LHC-b or high luminosity  $Z$  factory. As we have used generalised factorization approximation along with BSW model and Lepage and Brodsky model for penguin and annihilation contributions, we might have introduced certain uncertainties. Nevertheless, since the branching ratio obtained for  $B^- \rightarrow D^0 D_s^-$  matches very well with the experimental value, we point out here that the decay mode  $B^- \rightarrow D^0 D^-$  may also be pursued further in the first round of B-factory experiments (where it can easily be accessible) to observe direct CP violation or to provide us a hint for the presence of new physics.

## 5 Acknowledgements

AKG would like to thank Council of Scientific and Industrial Research, Government of India, for financial support.

## References

- [1] *CP Violation* ed. C. Jarlskog (World Scientific, Singapore, 1989)
- [2] *CP Violation* by I. I. Bigi and A. I. Sanda, Cambridge Monographs on Particle Physics, Nuclear Physics and Cosmology, (2000).
- [3] *CP Violation* by G. C. Branco, L. Lavoura and J. P. Silva, International series of Monographs on Physics, Number 103, Oxford University Press (1999).
- [4] N. Cabibbo, Phys. Rev. Lett. **10**, 531 (1963); M. Kobayashi and T. Maskawa, Prog. Theo. Phys. **49**, 652 (1973).

- [5] A. J. Buras and R. Fleischer, in *Heavy Flavours II* eds. A. J. Buras and M. Lindner (World Scientific, Singapore, 1998) p. 65.
- [6] *The BaBar Physics Book* eds. P. F. Harrison and H. R. Quinn (SLAC report 504, 1998).
- [7] L. L. Chau and W. -Y. Keung, Phys. Rev. Lett. **53**, 1802 (1984); C. Jarlskog and R. Stora, Phys. Lett. **B 208**, 268 (1988).
- [8] A. B. Carter and A. I. Sanda, Phys. Rev. Lett. **45**, 952 (1980); Phys. Rev. **D 23**, 1567 (1981); I. I. Bigi and A. I. Sanda, Nucl. Phys. **B 193**, 85 (1981).
- [9] A. Abashian *et al.* [BELLE Collab.], Phys. Rev. Lett **86**, 2509 (2001).
- [10] B. Aubert *et al.* [BaBar Collab.], Phys. Rev. Lett **86**, 2515 (2001).
- [11] T. Affolder *et al.* [CDF Collab.], Phys. Rev. **D 61**, 072005 (2000).
- [12] V. Fanti *et al.* [NA48 Collab.], Phys. Lett **B 465**, 335 (1999).
- [13] A. Alavi-Harati *et al.* [KTeV Collab.], Phys. Rev. Lett **83**, 22 (1999).
- [14] Particle data Group, Review of Particle Physics, D. E. Groom et al, Euro. Phys. J. **C 15**, 1 (2000).
- [15] M. Bander, D. Silverman and A. Soni, Phys. Rev. Lett. **43**, 242 (1979); J. M. Gérard and W. S. Hou, Phys. Rev. **D 43**, 2909 (1991).
- [16] A. I. Sanda and Z.-Z. Xing, Phys. Rev. **D 56**, 341 (1997).
- [17] Z.-Z. Xing, Phys. Lett. **B 443**, 365 (1998).
- [18] X.-Y. Pham and Z.-Z. Xing, Phys. Lett. **B 458**, 375 (1999).
- [19] Z.-Z. Xing, Phys. Rev. **D 61**, 014010 (2000).
- [20] D. Du and Z. Z. Xing, Phys. Lett. **B 349**, 215 (1995); D. Du, M. Z. Yang and D. Z. Zhang, Phys. Rev. **D 53**, 249 (1996).
- [21] M. Bauer, B. Stech and M. Wirbel, Zeit. Phys. **C 34**, 103 (1987).
- [22] G. P. Lepage and S. J. Brodsky, Phys. Lett. **B 87**, 359 (1979).
- [23] G. Buchalla, A.J. Buras and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
- [24] Y. H. Chen, H. Y. Cheng, B. Tseng and K. C. Yang, Phys. Rev. **D 60**, 094014 (1999); A. Ali, G. Kramer and C. D. Lü, Phys. Rev. **D 58**, 094009 (1998); A. Ali and C. Greub, Phys. Rev. **D**, 57, 2996 (1998).

Table 1: Numerical value of the effective Wilson coefficients  $c_i^{eff}$  for  $b \rightarrow s$  and  $b \rightarrow d$  transitions.

|                       | $b \rightarrow s$     |                       | $b \rightarrow d$     |                       |
|-----------------------|-----------------------|-----------------------|-----------------------|-----------------------|
|                       | Time-like             | Space-like            | Time-like             | Space-like            |
| $c_1^{eff}$           | 1.168                 | 1.168                 | 1.168                 | 1.168                 |
| $c_2^{eff}$           | -0.366                | -0.366                | -0.366                | -0.366                |
| $c_3^{eff}$           | 0.0225+i0.0044        | $-(0.0096 + i0.0003)$ | 0.0197+i 0.005        | $-(0.0123 - i0.0066)$ |
| $c_4^{eff}$           | $-(0.0456 + i0.0133)$ | $(0.0505 + i0.0009)$  | $-(0.0373 + 0.015)$   | $(0.0586 - i0.0199)$  |
| $c_5^{eff}$           | 0.0132+i0.0044        | $-(0.0189 + i0.0003)$ | 0.0104+i0.005         | $-(0.0216 - i0.0066)$ |
| $c_6^{eff}$           | $-(0.0478 + i0.0133)$ | $(0.0483 + i0.0009)$  | $-(0.0395 + i0.015)$  | $(0.0564 - i0.0199)$  |
| $c_7^{eff}/\alpha$    | $-(0.0282 + i0.0363)$ | $-(0.0303 - i0.0018)$ | $-(0.0119 + i0.0398)$ | $-(0.0143 + i0.0391)$ |
| $c_8^{eff}/\alpha$    | 0.055                 | 0.055                 | 0.055                 | 0.055                 |
| $c_9^{eff}/\alpha$    | $-(1.4252 + i0.0363)$ | $-(1.4273 - i0.0018)$ | $-(1.4089 + i0.0398)$ | $-(1.4113 + i0.0391)$ |
| $c_{10}^{eff}/\alpha$ | 0.481                 | 0.481                 | 0.481                 | 0.481                 |

- [25] A. J. Buras et al, Nucl. Phys. **B 370**, 69 (1992); M. Ciuchini et al, Zeit. Phys. **C 68**, 239 (1995).
- [26] Y. H. Chen, H. Y. Cheng, B. Tseng and K. C. Yang, Phys. Rev. **D 60**, 094014 (1999).
- [27] N. G. Deshpande and J. Trampetic, Phys. Rev. **D 41**, 2926 (1990).
- [28] H. Y. Cheng and K. C. Yang, Phys. Rev. **D 59**, 092004, (1999).
- [29] A. Khodjamirian and R. Rückl, in *Heavy Flavours II*, ed. A. J. Buras and M. Lindner, World Scientific Singapore (1998).
- [30] A. Ali, G. Kramer and C. D. Lü, Phys. Rev. **D 59**, 014005 (1998).

Table 2: Branching ratio and CP Asymmetry in % for  $B^- \rightarrow D^0 D_s^-, D^0 D^-$  decay modes.

| Decay modes                 | Branching             | Ratios                | $(BR)$<br>Expt.                | $CP$ Asymmetry  |                 |
|-----------------------------|-----------------------|-----------------------|--------------------------------|-----------------|-----------------|
|                             | $N_c^{eff} = 2$       | $N_c^{eff} = 3$       |                                | $N_c^{eff} = 2$ | $N_c^{eff} = 3$ |
| $B^- \rightarrow D^0 D_s^-$ | $1.29 \times 10^{-2}$ | $1.46 \times 10^{-2}$ | $(1.3 \pm 0.4) \times 10^{-2}$ | 0.18            | 0.18            |
| $B^- \rightarrow D^0 D^-$   | $8.72 \times 10^{-4}$ | $9.86 \times 10^{-4}$ | $< 6.7 \times 10^{-3}$         | 3.62            | 3.62            |